

# Chiral Dynamics beyond the Standard Model

Jan Stern<sup>a</sup>

<sup>a</sup>Theory Group, IPN-Orsay, Universite de Paris Sud XI, 91406 ORSAY, France

The SM Lagrangian without physical scalars is rewritten as the LO of a Low-Energy Effective Theory invariant under a higher non linear symmetry  $S_{nat} \supset SU(2)_W \times U(1)_Y$ . Soft breaking of  $S_{nat}$  defines a hierarchy of non standard effects dominated by universal couplings of right handed quarks to W. The interface of corresponding EW tests with non perturbative QCD aspects is briefly discussed.

## 1. ELECTROWEAK LOW-ENERGY EFFECTIVE THEORY

A systematic search of Physics beyond the Standard Model is usually formulated in the framework of a Low-Energy Effective theory (LEET). One expects that at very high energies  $E > \Lambda$  there exist new (gauge) particles and that their interaction is governed by new (local) symmetries not contained in the SM gauge group  $SU(2)_W \times U(1)_Y$ . In order to specify the LEET which would describe the Physics below the scale  $\Lambda$  in a systematic low-energy expansion, one has to identify the characteristic property which makes from the SM the unique and precise effective description of low-energy phenomena. The most popular approach is provided by the

### 1.1. Decoupling Scenario

One assumes that below the scale  $\Lambda$ , it is possible to ignore (integrate out) the heavy states and forget the corresponding (gauge) symmetries beyond  $SU(2)_W \times U(1)_Y$ , i.e. to remain with the degrees of freedom and symmetries of the SM. The latter is then identified as the most general  $SU(2)_W \times U(1)_Y$  invariant Lagrangian that can be constructed out of the SM fields and is **renormalizable order by order in powers of coupling constants**. The effective Lagrangian then reads

$$\mathcal{L}^{eff} = \mathcal{L}_{SM} + \sum_{D>4} \frac{1}{\Lambda^{D-4}} \mathcal{O}_D \quad (1)$$

where  $D$  denotes the mass (ultraviolet) dimension of the operator  $\mathcal{O}_D$ . Using the UV dimension  $D$  as the sole indication of the degree of suppression of a non standard operator is related to the emphasize put on renormalizability. However, it is not very practical: At the NLO order, there are 80 independent gauge invariant  $D=6$  operators [6] and a finer classification of new operators might be needed.

### 1.2. A not quite decoupling alternative

Even if below the scale  $\Lambda$  the heavy particles decouple, their interaction may be such that some high energy symmetries  $S_{nat}$  beyond  $SU(2)_W \times U(1)_Y$  survive at low energies and constrain the LEET. Since below the scale  $\Lambda$  the set of effective fields - typically the SM fields - is too small to span a linear representation of  $S_{nat}$ , the extra symmetry relevant at low energy

$$\frac{S_{nat}}{SU(2)_W \times U(1)_Y} \equiv \mathcal{C}_{sp} \quad (2)$$

should be realized nonlinearly. This means that the symmetry (2) is not manifest in the low-energy spectrum but it constrains interaction vertices. The effective Lagrangian should exhibit the symmetry  $S_{nat}$  and should be organized in a systematic low-energy expansion in powers of momenta

$$\mathcal{L}^{eff} = \sum_{d \geq 2} \mathcal{L}_d \quad (3)$$

where  $\mathcal{L}_d = \mathcal{O}(p^d)$  in the low-energy limit. (In general, the chiral dimension  $d \neq D$ . A concise

review of infrared power counting can be found in [3] , [5].) The resulting non decoupling LEET is **renormalized and unitarized order by order in the momentum expansion (3)**, following and extending the example of Chiral Perturbation Theory (ChPT), [1] [2]. Under these circumstances there is no particular reason to start the LE expansion around a renormalizable limit: In its minimal version, the LEET may contain all the observed particles of the SM without the physical Higgs scalar, unless such a light scalar is found experimentally. On the other hand, the presence of three GB fields  $\Sigma \in SU(2)$  is crucial to generate the masses of W and Z by the standard Higgs mechanism. The standard gauge boson mass term coincides with the Goldstone boson kinetic term

$$\mathcal{L}_{mass} = \frac{1}{4} F_W^2 \text{Tr}(D_\mu \Sigma^\dagger D^\mu \Sigma). \quad (4)$$

(This becomes manifest in the physical gauge  $\Sigma = 1$ ). An important difference between the decoupling and non decoupling LEET concerns the scale  $\Lambda$  above which the effective description breaks down. Whereas in the decoupling case  $\Lambda$  is essentially independent of the low-energy dynamics itself ( the latter is renormalizable) and it cannot be estimated apriori, the consistency of the loop expansion in the decoupling case requires

$$\Lambda \approx \Lambda_W = 4\pi F_W \approx 3TeV. \quad (5)$$

### 1.3. What is $S_{nat}$ ?

The experimental fact that at the leading order of the low-energy expansion (3) one meets the bare SM interaction vertices should now follow from the higher symmetry

$$S_{nat} \supset S_{EW} \equiv SU(2)_W \times U(1)_Y \quad (6)$$

and not from the requirement of renormalizability as in eq (1) above. Indeed, as shown in ref [4] , given the set of SM fields observed at low energy and their standard transformation properties under  $S_{EW}$ , one can construct several invariants of  $S_{EW}$  carrying the leading infrared dimension  $d = 2$  that are not present in the SM Lagrangian. Such operators are not observed and it is a primary role of the symmetry  $S_{nat}$  to suppress them. This presumes that  $S_{nat}$  is a non linearly

realized (hidden) symmetry of the SM itself and it should be possible to infer it from the known SM interaction vertices. This turns out to be the case [4] , [5], provided one sticks to the part  $\mathcal{L}'_{SM}$  of the SM Lagrangian that does not contain the physical Higgs scalar nor the Yukawa couplings to fermions. Proceeding by trial and error one can show that the condition  $\mathcal{L}_2 = \mathcal{L}'_{SM}$  admits a unique **minimal solution**

$$S_{nat} = [SU(2) \times SU(2)]^2 \times U(1)_{B-L}. \quad (7)$$

This provides a guide of constructing the LEET below the scale  $\Lambda_W$  and it might indicate which new heavy particles could be expected well above this scale.

### 1.4. Spurions

It is conceivable that at ultrahigh energies the symmetry  $S_{nat}$  is linearly realised via 4  $SU(2)$  and one  $U(1)$  gauge fields with 9 of them acquiring a mass  $> \Lambda_W$ . As the energy decreases below  $\Lambda_W$ , this linearly realized symmetry is reduced ending up with  $S_{EW} = SU(2)_W \times U(1)_Y$  spanned by the electroweak gauge bosons. This reduction may be viewed as a pairwise identification of two independent  $SU(2)$  factors in (7) **up to a gauge transformation**  $\Omega(x)$ . Denoting the corresponding  $SU(2)_I \times SU(2)_{II}$  connections by  $A_\mu^I$  and  $A_\mu^{II}$  respectively, the identification amounts to the constraint

$$A_\mu^I = \Omega A_\mu^{II} \Omega^{-1} + i\Omega \partial_\mu \Omega^{-1} \quad (8)$$

This constraint is covariant under the original symmetry  $SU(2)_I \times SU(2)_{II}$  provided the gauge transformation  $\Omega(x)$  is promoted to a field transforming with respect to the  $SU(2)_I \times SU(2)_{II}$  group as the bifundamental representation

$$\Omega(x) \rightarrow V_I(x) \Omega(x) V_{II}(x)^{-1}. \quad (9)$$

The field  $\Omega(x)$  ( more precisely its multiple) is a remnant of the reduction procedure and we call it **spurion**: It does not propagate since its covariant derivatives vanishes  $D_\mu \Omega(x) = 0$  by virtue of eq.(8). There exists a gauge in which each spurion reduces to a constant multiple of the unite matrix. After the reduction of the original symmetry  $S_{nat}$  to  $S_{EW}$ , one remains with the 4 SM

gauge fields and with 3  $SU(2)$  valued scalar spurions populating the coset space

$$\mathcal{C}_{sp} = \frac{S_{nat}}{S_{EW}} = [SU(2)]^3 \quad (10)$$

and transforming as bifundamental representations of the symmetry group  $S_{nat}$ , as illustrated by the above example, c.f.(9). We refer the reader to [5], where more details of this construction, quantum number matching and unicity are given and discussed.

### 1.5. Summary

Before the symmetry  $S_{nat}$  is reduced via spurions, the content of the LEET consists of two disconnected sectors: First **the elementary sector** governed by the gauge group  $[SU(2)_L \times SU(2)_R \times U(1)_{B-L}]_{el}$  that acts on a set of (elementary) chiral fermion doublets transforming as  $[1/2, 0; B-L]$  and  $[0, 1/2; B-L]$  respectively, (similarly as in LR symmetric models [7]). Second, **the composite sector** containing three GBs  $\Sigma \in SU(2)$  arising from the spontaneous breakdown of the “composite” symmetry  $[SU(2)_L \times SU(2)_R]_c$  similarly to what happens in QCD. The covariant constraints reducing  $S_{nat} \rightarrow S_{EW}$  first identify (up to a gauge) the elementary and the composite  $SU(2)_L$ . This gives rise to a single  $SU(2)_W$ , the one of the SM, and to a spurion  $\mathcal{X}$  transforming as  $[1/2, 1/2]$  with respect to the composite  $\times$  elementary  $SU(2)_L$ . Analogously, in the right handed sector, one identifies the two  $SU(2)_R$  groups (spurion  $\mathcal{Y}$ ) and finally, the resulting right isospin is identified with  $U(1)_{B-L}$  (embedded into  $SU(2)$ ) (spurion  $\mathcal{Z}$ ). What remains is  $U(1)_Y$  where the hypercharge  $Y/2 = I_R^3 + (B-L)/2$  as required by the SM. Notice that the spurion  $\mathcal{Y}$  admits a gauge invariant decomposition  $\mathcal{Y} = \mathcal{Y}_u + \mathcal{Y}_d$  making appear projectors of the up and the down components of right handed doublets.

Hence, the structure of the SM vertices leads to three  $SU(2)$  valued spurions which in the physical gauge reduce to three parameters  $\mathcal{X} \rightarrow \xi$ ,  $\mathcal{Y} \rightarrow \eta$  and  $\mathcal{Z} \rightarrow \zeta$ . They are external to the SM and they parametrize the small explicit breaking of the symmetry  $S_{nat}$ . The fact that the spurion  $\mathcal{Z}$  necessarily carries a non zero value of  $B-L$  means that the lepton number violation is an unavoidable byproduct of the above construction [5]: Its

strength is tuned by the parameter  $\zeta \ll \xi$ ,  $\eta \ll 1$  and it cannot be anticipated within the LEET alone.

## 2. LEADING EFFECTS BEYOND THE STANDARD MODEL

The whole construction of  $\mathcal{L}^{eff}$  can now be devised in three steps:

- One collects all invariants under  $S_{nat}$  made up from the corresponding 13 gauge fields, GB field  $\Sigma$ , chiral fermions and the three spurions.
- One orders them according to the increasing chiral dimension  $d$  and to the number of spurion insertions.
- One imposes the invariant constraints between gauge fields and spurions eliminating the redundant degrees of freedom (c.f. (8)).

By construction, the leading term with  $d = 2$  and no spurion insertion coincides with the higgsless vertices of the SM with massive W and Z and massless fermions. The genuine effects beyond SM are identified with terms explicitly containing spurions. The latter naturally appear in a hierarchical order given by the power of the (small) spurionic parameters  $\xi$ ,  $\eta$  (we can consistently disregard the tiny spurion  $\zeta$  and LNV processes.) Despite the present inability of the LEET formalism to anticipate the actual size of individual non standard effects, we might well be able to order these effects i.e., **to predict their relative importance**.

### 2.1. Fermion Masses and power counting

Spurions are needed to construct a  $S_{nat}$  invariant fermion mass term: At leading order such terms are proportional to the operators

$$\mathcal{L}_{fm} = \bar{\Psi}_L \mathcal{X}^\dagger \Sigma \mathcal{Y}_a \Psi_R \quad (11)$$

where  $a = u, d$ . The chiral dimension of these operators is  $d = 1$ , ( $1/2$  for each fermion field) and they are furthermore suppressed by the spurionic factor  $\xi\eta$ . The mass term of the heaviest fermion (the top quark) with the chirally protected mass

should count at low energies as  $\mathcal{O}(p^2)$ . This suggests the following counting rule for the spurion parameters

$$\xi\eta \sim m_{top}/\Lambda_W \sim \mathcal{O}(p) \quad (12)$$

attributing to both  $\xi$  and  $\eta$  the chiral dimension  $d = 1/2$ . This leaves space for the existence of much lighter fermions with the mass term containing additional powers of  $\xi$  and  $\eta$ .

On the other hand, the understanding of **the smallness of neutrino masses** within the present LEET framework represents an alternative to the see-saw mechanism. The Majorana mass term necessarily involves the  $\Delta L = 2$  spurion  $\mathcal{Z}$  which brings in a new (tiny) scale related to LNV. Yet, one has to find the symmetry that suppresses the neutrino Dirac mass of the type (11). In [4] it has been proposed to associate this latter suppression with the discrete reflection symmetry

$$\nu_R^i \rightarrow -\nu_R^i. \quad (13)$$

This is possible since at the leading order the right handed neutrino does not carry any gauge charge and the symmetry (13) does not prevent the right handed neutrino to develop its own (small) Majorana mass. This mechanism of suppression of neutrino masses has yet another consequence. It forbids **the charged right-handed leptonic currents** despite the fact that  $\nu_R$  remains light. This corollary will take its importance later.

## 2.2. The full NLO

Sofar, all terms had the total chiral dimension  $d = 2$ , including (in the case of the mass term) the spurion factor according to the counting rule (12). This is characteristic of the LO. There is no  $d = 3$  term not containing spurions. On the other hand there are two and only two operators quadratic in spurions with the total chiral dimension  $d = 3$ . They represent non standard (i.e. containing spurions)  $S_{nat}$  invariant couplings of fermions to gauge bosons. In the case of left handed fermions the unique such operator reads

$$\mathcal{O}_L = \bar{\Psi}_L \mathcal{X}^\dagger \Sigma \gamma^\mu D_\mu \Sigma^\dagger \mathcal{X} \Psi_L, \quad (14)$$

whereas in the right handed sector the corresponding operator has four components ( $a, b =$

$u, d$ ) that are separately invariant under  $S_{nat}$

$$\mathcal{O}_R^{a,b} = \bar{\Psi}_R \mathcal{Y}_a^\dagger \Sigma^\dagger \gamma^\mu D_\mu \Sigma \mathcal{Y}_b \Psi_R. \quad (15)$$

Both operators are suppressed by a quadratic spurion factor: In view of the rule (12), they count as  $\mathcal{O}(p^2 \xi^2) \sim \mathcal{O}(p^3)$  and  $\mathcal{O}(p^2 \eta^2) \sim \mathcal{O}(p^3)$ , respectively. It is remarkable that the operators (14) and (15) cannot be generated by loops. This follows from the Weinbergs power counting formula

$$d = 2 + 2L + \sum_v (d_v - 2) \quad (16)$$

originally established in the case of ChPT [1] and subsequently extended to the case of LEET containing massive vector particles, chiral fermions [3] as well as spurions [4]. Eq (16) represents the chiral dimension of a connected Feynman diagram containing  $L$  loops and the vertices  $v$  of a chiral dimension  $d_v \geq 2$ . In the low-energy expansion the tree diagrams with a single insertion of a vertex (14) or (15) are more important than any loop contribution or than higher order purely bosonic trees. Accordingly, the NLO operators (14) and (15) are universal in the flavour space. Loops, oblique corrections, flavour dependence (and related FCNC), only come at the NNLO. The two operators (14) and (15) thus describe the potentially most important effects beyond the SM, **predicted** in the present framework of a non decoupling LEET. Unfortunately, it is not straightforward to translate this qualitative prediction into a more quantitative statement (see [8], [11]). The operators (14) and (15) both contain couplings of fermions to  $W$  and to  $Z$ . The latter involves too many apriori unknown LECs, especially in the right handed sector (15). For this reason, we first consider

## 2.3. Couplings of quarks to W at NLO

In the physical gauge and in a flavour basis in which both the mass matrix of  $u$  and  $d$  quarks are diagonal, the  $LO + NLO$  of a universal charged current coupled to  $W$  as defined by eqs (14) and (15) reads

$$\mathcal{L}_{CC} = g[l_\mu + \frac{1}{2}\bar{U}(\mathcal{V}_{eff}\gamma_\mu + \mathcal{A}_{eff}\gamma_\mu\gamma_5)D]W^\mu + hc, \quad (17)$$

where  $l_\mu$  stands for the standard leptonic charged  $V - A$  current not affected at NLO. The matrix notation is used in the flavour (family) space:  $U^T = (u, c, t)$ ,  $D^T = (d, s, b)$ , whereas  $\mathcal{V}_{eff}$  and  $\mathcal{A}_{eff}$  are complex  $3 \times 3$  effective EW coupling matrices. At NLO they take the form

$$\mathcal{V}_{eff}^{ij} = (1 + \delta)V_L^{ij} + \epsilon V_R^{ij} \quad (18)$$

$$\mathcal{A}_{eff}^{ij} = -(1 + \delta)V_L^{ij} + \epsilon V_R^{ij}, \quad (19)$$

where  $V_L = L_u L_d^\dagger$  and  $V_R = R_u R_d^\dagger$  with  $L_u, R_u$  and  $L_d, R_d$  denoting the pairs of unitary matrices that diagonalize the masses of the u-type and d-type quarks respectively. The parameters  $\delta = \mathcal{O}(\xi^2)$  and  $\epsilon = \mathcal{O}(\eta^2)$  originate from the spurions. Their magnitude is estimated to be of at most a fraction of per cent [5]. Hence at NLO, the LEET predicts two major non standard effects concerning the coupling of quarks to W:

i) The existence of **direct couplings of right handed quarks to W** (keeping in mind that the discrete symmetry (15) forbids similar couplings for leptons.)

ii) The chiral generalization of CKM unitarity and mixing.

At the LO one has  $\delta = \epsilon = 0$  and one recovers the CKM unitarity of the SM:  $\mathcal{V}_{eff} = -\mathcal{A}_{eff} = V_{CKM}$ . At the NLO, the two distinct matrices  $V_L$  and  $V_R$  are both unitary (for the same reason as in the SM) but the effective vector and axial-vector matrices  $\mathcal{V}$  and  $\mathcal{A}$  which are accessible in semi-leptonic transitions are not unitary anymore.

### 3. INTERFACE OF THE EW AND QCD EFFECTIVE COUPLINGS

A measurement of effective EW couplings  $\mathcal{V}_{eff}^{ij}$  and  $\mathcal{A}_{eff}^{ij}$  requires an independent knowledge of the involved non perturbative QCD parameters like the decay constants  $F_\pi, F_K, F_D, F_B$  or the transition form factors such as  $f_+^{K_0\pi^-}(0) \dots$ . The longstanding problem of an accurate extraction of the CKM matrix element  $V^{us}$  and the related test of the “CKM unitarity” illustrates this point and it becomes even more acute in the presence of non standard EW couplings, such as RHCs. The unfortunate circumstance is that the

most accurate experimental information on QCD quantities mentioned above, in turn come from semi leptonic transitions of the type  $P \rightarrow l\nu$  and  $P' \rightarrow Pl\nu$  where  $P = \pi, K, D, B$  and, consequently, the result of their measurement depends on (apriori unknown) EW couplings (18), (19). Finding an exit from this **circular trap** is a major task of phenomenological Flavor Physics.

#### 3.1. Non Standard EW Parameters in the Light Quark Sector

Let us first concentrate on light quarks  $u, d$ , and  $s$ . For them the SM loop effects simulating RHCs are strongly suppressed by at least two powers of mass.

Since  $V_L^{ub}$  is negligible and

$$\mathcal{V}_{eff}^{ud} = 0.97377(26) \equiv \cos\hat{\theta} \quad (20)$$

is very precisely known [9] from nuclear  $0^+ \rightarrow 0^+$  transitions, the light quark effective couplings  $\mathcal{V}_{eff}^{ua}, \mathcal{A}_{eff}^{ua}$ ,  $a = d, s$  can be expressed in terms of three non standard effective EW parameters: the spurion parameter  $\delta$  defined in (20), (21) and two RHCs parameters  $\epsilon_{NS}$  and  $\epsilon_S$  defined as

$$\epsilon V_R^{ud} = \epsilon_{NS} \cos\theta, \epsilon V_R^{us} = \epsilon_S \sin\theta, \quad (21)$$

where, upon neglecting  $V_L^{ub}$ , we have denoted

$$V_L^{ud} = \cos\theta = \cos\hat{\theta}(1 - \delta - \epsilon_{NS}), V_L^{us} = \sin\theta(22)$$

In the limit of SM all spurion NS parameters vanish,  $\delta = \epsilon_{NS} = \epsilon_S = 0$  and all light quark EW effective couplings are fixed by the experimental value of  $\mathcal{V}_{eff}^{ud}$ .

#### 3.2. $F_\pi, F_K, f_+(0)$ and $\mathcal{V}_{eff}^{us}$

Here  $F_\pi, F_K$  and  $f_+^{K^0}(0)$  stand for radiatively corrected genuine QCD quantities defined as residues of GB poles in two-point and three-point functions of axial and vector currents. These are the quantities that are subject to ChPT and/or lattice studies. It is further understood that all isospin breaking effects due to  $m_d - m_u$  are included. From the experimentally well known branching ratios [10]  $K(l2(\gamma))/\pi(l2(\gamma))$  one can infer  $\mathcal{A}^{us}F_K/\mathcal{A}^{ud}F_\pi$  and from the rate of  $K_{l3}^0$  we can extract the value of  $|f_+^{K^0}(0)\mathcal{V}^{us}|$ . Assuming the Standard Model couplings of quarks to

W this is sufficient to extract very accurate values that the corresponding QCD quantities would take in a world with vanishing spurion parameters  $\delta, \epsilon_{NS}, \epsilon_S$ . The latter are denoted by a hat. They read

$$\hat{F}_\pi = (92.4 \pm 0.2) \text{ MeV}, \hat{F}_K / \hat{F}_\pi = 1.182 \pm 0.007, \quad (23)$$

$$\hat{f}_+^{K^0}(0) = 0.951 \pm 0.005 \quad (24)$$

Here the errors merely reflect the experimental uncertainties in the measured branching ratios. In the presence of NS couplings of quarks to W the values of genuine QCD quantities extracted from semileptonic BRs are modified. Neglecting higher powers of spurion parameters  $\delta$  and  $\epsilon$  one gets using Eqs (21) and (22)

$$F_\pi^2 = \hat{F}_\pi^2 (1 + 4\epsilon_{NS}) \quad (25)$$

$$\left(\frac{F_K}{F_\pi}\right)^2 = \left(\frac{\hat{F}_K}{\hat{F}_\pi}\right)^2 \frac{1 + 2(\epsilon_S - \epsilon_{NS})}{1 + \frac{2}{\sin^2 \theta}(\delta + \epsilon_{NS})} \quad (26)$$

$$\left[f_+^{K^0\pi^-}(0)\right]^2 = \left[\hat{f}_+^{K^0\pi^-}(0)\right]^2 \frac{1 - 2(\epsilon_S - \epsilon_{NS})}{1 + \frac{2}{\sin^2 \theta}(\delta + \epsilon_{NS})} \quad (27)$$

The 3 NLO EW parameters  $\delta$ ,  $\epsilon_{NS}$  and  $\epsilon_S$  can be constrained using independent informations on QCD quantities  $F_\pi$ ,  $F_K$  or  $f_+^{K^0\pi^-}$ . Such information can originate from lattice simulations or from ChPT based measurements. As a matter of example let us mention the possible determination of  $F_\pi$  from the  $\pi_0 \rightarrow 2\gamma$  partial width or from precision  $\pi\pi$  scattering experiments, which are independent from the standard determination based on the  $\pi l 2$  decay rate. Despite a present lack of accuracy, this way may eventually provide a measurement of  $\epsilon_{NS}$  through Eqs (23) and (25).

The non standard EW parameters allow to infer the NLO deviation from the unitarity of the effective mixing matrix  $\mathcal{V}_{eff}$ , which is relevant in the description of  $K_{l3}$  decays. One finds

$$[\mathcal{V}_{eff}^{ud}]^2 + [\mathcal{V}_{eff}^{us}]^2 = 1 + 2(\delta + \epsilon_{NS} \cos^2 \hat{\theta} + \epsilon_S \sin^2 \hat{\theta}) \quad (28)$$

Due to the presence of RHCs, flavour mixing effects in  $F_K/F_\pi$  and in  $f_+^{K^0\pi^-}$  are no more related as in the SM.

### 3.3. Enhancement of $\epsilon_S$ ?

As already pointed out the genuine spurion parameters  $\delta$  and  $\epsilon$  can hardly exceed 0.01. On the other hand, the unitarity of the right-handed mixing matrix  $V_R$  implies

$$|\epsilon_{NS}|^2 \cos^2 \theta + |\epsilon_S|^2 \sin^2 \theta \leq \epsilon^2. \quad (29)$$

Since we live close to the left-handed world, one has  $\sin \theta \sim 0.22$ , reflecting the well known hierarchy of left handed flavour mixing. This in turn implies that  $|\epsilon_{NS}| < \epsilon$  remains tiny. On the other hand,  $|\epsilon_S| \leq 4.5\epsilon$  and it can indeed be enhanced to a few percent level, **provided the hierarchy in right-handed flavour mixing is inverted**, i.e.  $|V_R^{ud}| \ll |V_R^{us}|$ . In [8] it has been shown that a stringent test involving the EW coupling  $\epsilon_S - \epsilon_{NS}$  can be devised in  $K_{\mu 3}^L$  decay, (see [11]). On the other hand, it seems rather difficult to find another clean manifestation of RHCs driven by  $\epsilon_S$  and the remaining two NLO EW constants  $\delta$  and  $\epsilon_{NS}$  are presumably too small to be reliably detected. These remarks still hold if the preceeding discussion is extended to processes involving the short distance rather than chiral QCD dynamics.

## REFERENCES

1. S. Weinberg, Physica A96, 327, (1979).
2. J. Gasser and H. Leutwyler, Nucl. Phys. B250, 465 (1985).
3. J. Wudka, Int. J. Mod. Phys. A9, 2301, (1994).
4. J. Hirn and J. Stern, Eur. Phys. J. C 34 (2004), 447.
5. J. Hirn and J. Stern, Phys. Rev. D73, (2006), 056001.
6. W. Buchmueller and D. Wyler, Nucl. Phys. B268, (1986), 621.
7. R. N. Mohapatra, J.C. Pati, Phys. Rev. D11, (1975), 2558; G. Senjanovic, R. N. Mohapatra, Phys. Rev. D12 (1975), 1502.
8. V. Bernard, M. Oertel, E. Passemar and J. Stern, Phys. Lett B638, (2006), 480.
9. W. J. Marciano and A. Sirlin, Phys. Rev. Lett. 96, (2006), 032002.
10. F. Ambrosino et al (KLOE), Phys. Lett. B632, (2006), 43
11. M. Oertel, these proceedings.